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Two-dimensional superconductor-insulator transition in bulk single-crystal $YBa_2Cu_3O_{6.38}$

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We use a magnetic field to tune a highly anisotropic single crystal of oxygen-deficient $YBa_2Cu_3O_{7-\delta}$ with a transition temperature of 2 K through the superconductor-insulator transition. The sheet resistance scales with temperature, $0.05 \leq T \leq 1.0$ K, and field, $0 \leq H \leq 94$ kOe, in a manner predicted by a theory for quantum phase transitions in disordered two-dimensional superconductors.

The deleterious effects of disorder are amplified in lower dimensions, where long-range coherence is, at best, only marginally stable. A case in point is the destruction of superconductivity in ultrathin granular[1] and continuous[2] metal films. The disorder is parametrized by the film sheet resistance and the data suggest a zero-temperature threshold for quenching global superconducting coherence near $h/4e^2 = 6.45$ kΩ/□. Experimental variables for adjusting the disorder include film thickness for elemental metals[1-3] and composites,[4,5] stoichiometry or doping for compounds,[6,7] including oxide superconductor single crystals,[8] and ion bombardment.[9]

The application of a magnetic field to a disordered superconducting film leads to a fundamentally different type of $T=0$ superconductor-insulator transition.[10] Field-induced vortices Bose condense, resulting in a transition from a state with zero resistance and pinned vortices ("vortex glass") to a state with zero conductivity and localized Cooper pairs ("electron glass"). The finite-temperature scaling predictions of the Fisher theory[10] have been verified in amorphous-composite InO_x thin films,[11] and values for universal critical exponents predicted, but not numerically determined by the theory were measured. We report here evidence that the field-induced two-dimensional superconductor-insulator transition persists in a highly anisotropic three-dimensional superconductor, namely oxygen-deficient $YBa_2Cu_3O_{7-\delta}$ ($\delta = 0.62$). We find the same dynamical exponent describing the critical behavior of the sheet resistance in field as measured for α-InO_x, but with a very different scale for the transition. Our results on barely superconducting $YBa_2Cu_3O_{7-\delta}$ crystals permit a test of the field-induced superconductor-insulator scenario in a regime not accessible to disordered films as well as helping elucidate the nature of the incipient superconducting state in relatively pristine, but essentially decoupled CuO_2 planes.

Decreasing the oxygen content of single-crystal $YBa_2Cu_3O_{7-\delta}$ preferentially creates vacancies in the chains, reducing the superconducting-transition temperature T_c and increasing the anisotropy.[12] Moreover, when thermally equilibrated and quenched from elevated temperatures T_Q below 200 °C, the vacancies disorder, lowering T_c while keeping the total oxygen content of the crystal constant.[13] By adjusting T_Q, T_c can be precisely controlled. This unique property of $YBa_2Cu_3O_{7-\delta}$ allowed us to fine tune T_c sufficiently close to the superconductor-insulator boundary that a magnetic field could drive the transition.

We started with a high-quality twinned single crystal of $YBa_2Cu_3O_{7-\delta}$ with a superconducting-transition temperature of 13 K and a 10%–90% width of 1 K as measured in a SQUID magnetometer. The sample was heated to $T_Q = 150$ °C, quenched by immersion in liquid nitrogen, mounted on the sample holder and top loaded into a helium dilution refrigerator, resulting in a resistive-zero $T_{c0} = 2.1$ K. Four-probe magnetoresistance measurements in the a-b plane with $\mathbf{H} \parallel \hat{\mathbf{c}}$ were made for $0.05 \leq T \leq 15$ K and $0 \leq H \leq 94$ kOe using a standard lock-in technique at 16 Hz. Input power was restricted to less than 10^{-10} watts to avoid sample heating and nonlinear I-V characteristics. Two opposite faces of the sample were covered with silver epoxy to serve as low-resistance current leads and two 75-μm stripes were painted along $\hat{\mathbf{c}}$ on one adjoining face to serve as the voltage leads. A heat treatment at 550 °C served both to adjust the oxygen stoichiometry of the crystal and to sinter the silver to the sample surface, providing reliable low-resistance contacts (<1 Ω at $T = 0.05$ K in field). Effective sample dimensions were $0.6 \times 0.6 \times 0.3$($c$ axis) mm^3. Errors due to finite contact size and variations in the sample thickness make absolute values of $R(H,T)$ uncertain within 10%. Uncertainties due to possible microcracking from quenching the crystal have not been included. As the resistive anisotropy of barely superconducting $YBa_2Cu_3O_{7-\delta}$ is in excess of 3000,[12] the method used here of effectively shorting out any possible c-axis transport is crucial for obtaining the true resistance in the a-b plane.

We show in Fig. 1 $R_{sq}(T)$ in a semilogarithmic plot for selected values of H. The measured three-dimensional sample resistivity was converted to a two-dimensional resistance R_{sq} of *pairs* of CuO_2 planes by dividing the thickness of the crystal by the c-axis lattice constant,[14] 11.8 Å. The decision to use pairs of CuO_2 planes as the two-dimensional unit is based on recent measurements[15] of the Hall effect in fully oxygenated $YBa_2Cu_3O_{7-\delta}$ which conclude that the pairs of CuO_2 planes in each unit cell remain strongly coupled near T_c. The alternative assumption that individual CuO_2 planes are the noninteract-

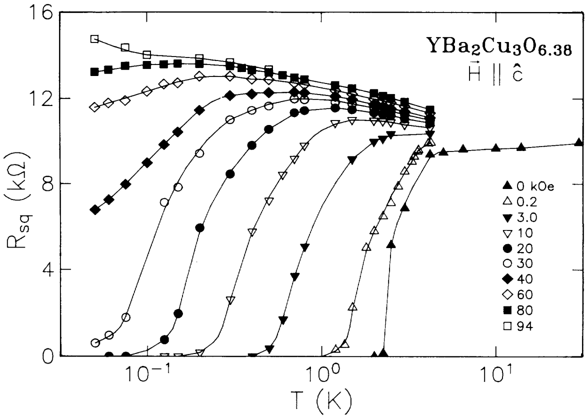

FIG. 1. Sheet resistance R_{sq} vs temperature T for a series of magnetic fields H perpendicular to the CuO_2 planes.

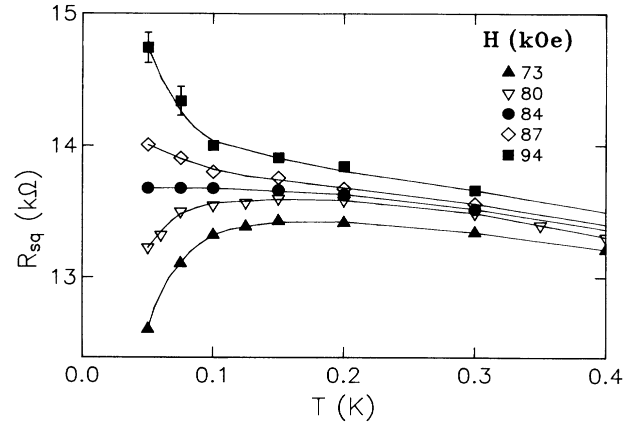

FIG. 2. Temperature dependence of the sheet resistance as $T \rightarrow 0$ in the immediate vicinity of the superconductor-insulator transition at $H_c = 84$ kOe and $R_{sq} \sim h/2e^2$.

ing two-dimensional unit has the consequence of doubling all values of R_{sq} reported here, but leaves our conclusions unchanged. At all T, $R_{sq} \propto H$ for $H > 70$ kOe. For $T \leq 0.075$ K and $H > 80$ kOe, we were unable to make equilibrium measurements of R_{sq} due to long relaxation times. Whether these long times are related to the glassy nature of the system at the transition or whether they reflect inadequate sample heat sinking at highest H and lowest T is unclear. In any case, we used the linear dependence of R_{sq} on field for 70 kOe $\leq H \leq$ 80 kOe to extrapolate to $H > 80$ kOe at $T = 0.05$ and 0.075 K. The identical extrapolation procedure gave values which agreed within 1% with the measured R_{sq} ($H = 94$ kOe) for all $T \geq 0.1$ K.

There are several features of interest in Fig. 1. First, the temperature dependence of the $H=0$ normal-state resistance is metalliclike above $T_{c0} = 2.1$ K. Given the large values of resistive anisotropies observed in this system and the fact that $\rho_c(T)$ is typically strongly upturned, this indicates that the contact configuration employed has negated any contributions from the c-axis resistivity. Second, the values of R_{sq}, whether in the normal or the mixed state, are of order $h/4e^2$ as found for two-dimensional superconductors near the disorder-induced superconductor-insulator transition.[16] Finally, as the low-temperature resistance at constant field proceeds from being monotonically decreasing to monotonically increasing with decreasing temperature, there must be a critical field H_c where R_{sq} is temperature independent.

We focus in Fig. 2 on the magnetic-field region about H_c in a linear plot of R_{sq} vs T. A field of 84 kOe causes the $T \rightarrow 0$ resistance to be constant, with $R_{sq} = 13.6$ k$\Omega \sim h/2e^2$. Within the context of the scaling theory of the field-induced superconductor-insulator transition,[10] this identifies $H_c = 84$ kOe as the critical field separating the "true" superconducting state with a pinned-vortex glass at $T=0$ from an insulating electron glass with Bose condensed vortices at $T=0$. In contrast to the pure three-dimensional case, the resistivity in the two-dimensional insulating glass is expected to increase with decreasing temperature slower than exponentially, consistent with the weak upturn which we observe. Furthermore, the theory predicts a scaling form,

$$R_{sq}(H,T) = (h/4e^2)\tilde{R}[c_0|H - H_c|/T^{1/z_B\nu_B}] ,$$

where $z_B \nu_B$ is a universal dynamical exponent and c_0 is a nonuniversal constant. Since

$$\left.\frac{dR_{sq}}{dH}\right|_{H_c} = \left(\frac{h}{4e^2}\right) c_0 T^{-1/z_B\nu_B} \tilde{R}'(0) , \qquad (1)$$

the product of the universal exponents $z_B \nu_B$ can be determined directly from the experimental results for the temperature dependence of (dR_{sq}/dH) evaluated at H_c.

We plot in Fig. 3 $(dR_{sq}/dH)|H_c$ vs T on a log-log scale. For $T < 1$ K the data follow a power-law form, and we find from Eq. (1) that $z_B \nu_B = 1.37 \pm 0.1$. This value for the dynamical exponent agrees well with the results[11] for

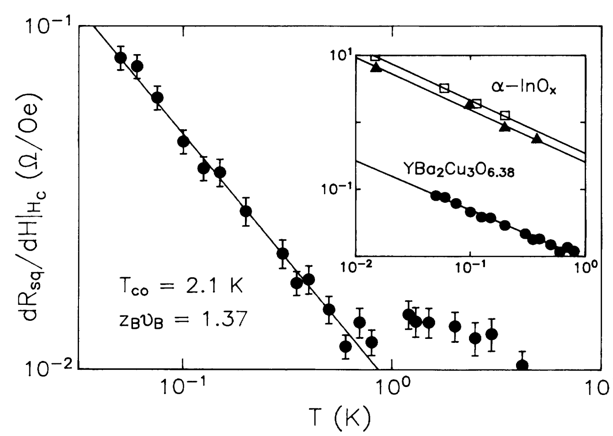

FIG. 3. Scaling behavior of the magnetic-field derivative of the sheet-resistance evaluated at H_c for $T < 1$ K, giving a dynamical exponent $z_B \nu_B = 1.37 \pm 0.1$. The solid line is a least-squares fit. The inset compares the scaling forms of single-crystal $YBa_2Cu_3O_{6.38}$ and amorphous-composite InO_x films (Ref. 11), with similar exponents but dissimilar prefactors.

amorphous-composite InO_x in field, where $z_B \nu_B = 1.26$ for a $T_c = 0.29$-K film and $z_B \nu_B = 1.31$ for a $T_c = 0.54$-K film. The experimental values for the exponent product are consistent with the theoretical constraints[10] $z_B = 1$ and $\nu_B \geq 1$.

We show in the inset to Fig. 3 the correspondence between the $YBa_2Cu_3O_{6.38}$ and the α-InO_x data sets. Although the exponents are essentially the same for the two systems, dR_{sq}/dH evaluated at H_c is of order 1 Ω/Oe for InO_x, while it is of order 0.03 Ω/Oe for $YBa_2Cu_3O_{6.38}$. This difference could be due to either the fact that single-crystal $YBa_2Cu_3O_{6.38}$, with disorder primarily in the chains, is closer to the clean limit or that c_0 is sensitive to possible weak coupling between the CuO_2 planes. In the latter view, c_0 gradually vanishing would be a natural way to pass from the two-dimensional to the three-dimensional limit.

Hebard and Paalanen[11] were able to test the prediction $H_c \propto T_c^{2/z}$ for a series of α-InO_x films, finding $z = 1$ as predicted. Unfortunately, we were unable to quench our crystal to lower T_c and make an independent test of the variation of H_c with T_c. Hence, the scaling behavior of Fig. 3 with a universal value for $z_B \nu_B$ remains the strongest evidence that bulk single crystals of $YBa_2Cu_3O_{7-\delta}$ exhibit a two-dimensional field-induced superconductor-insulator transition. We are able to compare, however, our result with the H_c vs T_c curve determined for α-InO_x. For $H_c = 84$ kOe one would expect a Kosterlitz-Thouless[17] transition temperature $T_c = 1.3$ K, which is close to the temperature where we observe the onset of scaling behavior. The disappearance of scaling behavior for $T > 1$ K may be due to finite-temperature corrections which are of order $(T/T_c)^2$ or it may reflect a crossover from two-dimensional to three-dimensional character as T approaches T_{c0}. Alternatively, the finite width of the superconducting transition may suppress artificially the true scaling form.

Previous studies[18] of the irreversibility line in oxygen-deficient single-crystal $YBa_2Cu_3O_{7-\delta}$ found large flux flow even for superconducting-transition temperatures as low as 7 K. We have shown here that the magnetic field-induced superconductor-insulator transition for single-crystal $YBa_2Cu_3O_{6.38}$ with $T_{c0} = 2.1$ K is dominated by the Bose condensation of the magnetic vortices at $H_c = 84$ kOe and $R_{sq} \sim h/2e^2$. The scaling behavior of $R_{sq}(H,T)$ at the transition is consistent with a two-dimensional description, establishing anisotropic oxide superconductors as fertile ground for testing universal behavior at the superconductor-insulator boundary.

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